# SUMMARY OF THE INTERNATIONAL SYMPOSIUM ON HEAVY FLAVOR AND ELECTROWEAK THEORY

R. D. Peccei
*Department of Physics & Astronomy, University of California*
*Los Angeles, CA 90095-1547, USA*
E-mail: peccei@physics.ucla.edu

ABSTRACT

This summary discusses some of the topics which were of overarching interest at the Symposium. These included, corrections to perturbative QCD predictions; heavy quark physics; electroweak symmetry breaking; and physics of the top quark.

## 1. Introductory Remarks.

Even with the best of will, it is impossible to summarize in 40 minutes the 30 talks given at the Symposium. Thus, instead I will try to concentrate on a few topics of overarching interest. These included, corrrections to perturbative QCD predictions; heavy quark physics; electroweak symmetry breaking; and physics of the top quark. There were many other interesting topics discussed at the Symposium [perturbation theory resummation; renormalons; CP and automorphisms; mass shifts in strong magnetic fields; symmetry pattern of mass matrices; etc.] which I, unfortunately, cannot properly cover in this summary. I apologize for this and refer the interested reader to the appropriate contributions in these Proceedings.

## 2. Corrections to Perturbative QCD.

One of the recurring themes in the Symposium was that perturbative QCD has its limitations. Perturbative QCD gives accurate predictions as long as the expansion parameter for the process in question is really $\alpha_s$. However, when this is not really the case, to obtain reliable predictions, one must include corrections which depend in detail on the physics of the problem. Three examples were discussed at the Symposium, each of which illustrated a particular way in which the relevant physics dictated how to augment the perturbative QCD calculations. Brodsky[1] considered threshold effects in heavy quark production in $e^+e^-$ collisions; Berger[2] discussed resumming initial state bremsstrahlung in top production at hadronic machines; and Wise[3] explained the role that color octet contributions have in hadronic production of charmonia. In each of these examples the underlying physics which causes modifications to perturbative QCD is quite clear. Indeed, for the processes discussed by Brodsky and Berger analogous phenomena occur also in QED. Nevertheless, each of these examples is a challenging area for QCD, if one wants accurate predictions to

compare with experiment.

Threshold production of pairs of charged fermions is sensistive to Coulomb exchange. For $e^+e^- \to \tau^+\tau^-$ near $\tau$-threshold one must include the multiple Coulomb rescattering of the produced pairs. Similarly, for heavy quark-antiquark production for $\beta = \sqrt{1-4m_Q^2/s} \to 0$ one must take into account of the gluonic Coulomb rescattering. For both QED and QCD one incorporates these effects through the introduction of a Coulomb factor, which sums up the multiple exchanges of photons or gluons:

$$S(x) = \frac{x}{1-e^{-x}}$$

with

$$x = \begin{cases} \frac{\pi}{\beta}\alpha & \text{QED} \\ \frac{\pi}{\beta}\frac{4}{3}\alpha_s & \text{QCD} \end{cases}$$

This Coulomb factor modifies the angular distribution at threshold, so that the coefficient of $\cos^2\theta$ is not simply $\beta^2$ but $\beta^2 S(x)$:

$$\frac{d\sigma}{d\Omega} \sim [2-\beta^2+\beta^2 S(x)\cos^2\theta] \ .$$

What Brodsky points out is that when one does the summing of the Coulomb exchanges at threshold properly, one obtains in this Coulomb factor the running coupling responsible for the binding of quarkonia $\alpha_V$, since the same physics is involved. Thus for QCD really one has

$$x = \frac{4\pi}{3}\frac{\alpha_V(\beta^2 s)}{\beta} \ .$$

This being the case, it may be possible to extract the coupling responsible for the charmonium bound state spectrum by studying the threshold angular distribution for $e^+e^- \to c\bar{c}$. A real question, however, is if this angular distribution is reflected faithfully in the angular distribution of the corresponding charmed hadrons, or whether hadronization effects mask entirely the Coulomb rescattering physics.

Berger[2] discussed another example where to properly calculate the physics of the problem one again has to sum up the effects of soft gluons–in his case, radiated gluons form the initial state. At the Tevatron the production of top quarks comes dominantly from the process $q\bar{q} \to t\bar{t}$. In contrast, at the LHC this will occur mostly through gluon fusion. In the usual fashion, the hadronic cross section for top production is then given by the convolution of the parton cross section and the quark and antiquark distribution functions

$$\sigma_{t\bar{t}}(s) = \int dx_1 dx_2 \ q(x_1)\bar{q}(x_2) \ \hat{\sigma}_{t\bar{t}}(x_1 x_2 \hat{s})$$

The partonic cross section $\hat{\sigma}_{t\bar{t}}$ is known to $O(\alpha_s^2)$. However, near threshold there are large corrections arising from the bremsstrahlung of a soft gluon ($p_g \to 0$) from the

initial state quarks or antiquarks. The single bremsstrahlung of a gluon introduces a factor

$$sb = \int_0^1 dz[1 + 2\alpha_s \ln(1-z)]$$

which, although finite, is large due to the soft gluon contribution at $z \to 1$ ($p_g \to 0$ corresponds to $z \to 1$). Thus, one should really consider also multiple soft gluon emission. As Berger discusses, one can actually resum the bremsstrahlung logarithms $(\alpha_s \ln(1-z))^n$ from multiple gluon emission and eventually one obtains a full enhancement factor of the form[4]

$$E \sim \alpha_s((1-z)^{2/3}m^2)\ln^2(1-z) \; .$$

However, from the above formula one sees that as $z \to 1$ one gets into scale values of $\alpha_s$ which are no longer in the perturbative regime.

There are different approaches of how to handle this. For instance, in this Symposium Berger[2] discussed how one can use a principal value regularization prescription to estimate the infrared sensitive part of the enhancement factor. However, the important message is that, because of these threshold effects, there is a bit of the top cross section at the Tevatron that is **uncalculable** in perturbative QCD. In fact, as Berger reported, what he and Contapanagos[5] do is to effectively set the resummed contribution to zero for $\eta < 0.005$ in the partonic cross section because they cannot trust the answer below this value. They obtain in this way for the top cross section at $\sqrt{s} = 1.8$ TeV, assuming $m_t = 175$ GeV, the value

$$\sigma_{t\bar{t}}(1.8 \text{ TeV}) = (5.5 \pm 0.3) \text{ pb} \; .$$

Here the error is an estimate of the uncertainty coming from the structure functions and the scale uncertainties. Because the resummed contribution contributes about 0.5 pb to the top cross section, the error coming from the excluded region near $\eta = 0$ probably is not significant. Nevertheless, it would be nice to have an estimate also of its possible magnitude.

Wise[3] discussed some aspects of charmonium production in hadronic collisions. This is a topic of considerable interest since recent data at the Tevatron showed that the production of $\psi, \psi'$ and $\Upsilon$ is much larger than was expected from a perturbative QCD quarkonia calculation[6]. Schematically, quarkonium production is given by convoluting the partonic cross-section for producing gluons of a certain fractional momentum with the gluon fragmentation function for quarkonia:

$$d\sigma(p) = \int dz \; \hat{\sigma}(z) P(p/z) D_{g\to Q\bar{Q}}(z)$$

A naive estimate of the gluon fragmentation function can be obtained by considering the same graphs which contribute to quarkonium decay. This gives for states whose

decay involve two gluons

$$D_{g\to Q\bar{Q}}(z) \sim \frac{\alpha_s^2}{\pi}\frac{|\psi|^2}{m_Q^2} f(z) \sim \alpha_s^2 v^{3+2L} \ .$$

Here $v$ is the relative velocity of the bound quarks and $L$ is the angular momentum associated with the produced quarkonia.

In his talk, Wise[3] emphasized that because one is dealing with bound state production one cannot just naively apply the same ideas that hold in quarkonium decay. Thus, for example, for the $L = 1$ $\chi$-states besides the naive result for $D_{g\to\chi} \sim \alpha_s^2 v^5$, one can imagine[7] also production via an $L = 0$ color octet intermediate state which then decays via soft gluon emission to the $\chi$. Such a color octet contribution still involves a factor of $v^5$ but now is proportional to $\alpha_s$ not $\alpha_s^2$

$$D^8_{g\to\chi} \sim \alpha_s v^5$$

and hence, in principle, can give a much larger contribution. Similar considerations hold for $\psi$ production, where the naive quarkonium estimate gives for the production of the $L = 0$ $c\bar{c}$ state

$$D_{g\to\psi} \sim \alpha_s^3 v^3 \ ,$$

while the contribution arising from an $L = 0$ color octet intermediate state, which then decays into a $\psi$ by emitting two soft gluons[7], gives

$$D^8_{g\to\psi} \sim \alpha_s v^7 \ .$$

One gains a factor of $\alpha_s^2$ but at the price of a $v^4$ factor. So here it is not so clear whether the color octet contribution can give an enhancement.

Because detailed bound state calculations are not simple to do, it is difficult to estimate reliably how much each of the above mechanisms really contributes to the gluon fragmentation function into quarkonia. Thus, it might be very useful to have a diagnostic test which may help distinguish among these different mechanism. Wise[3] suggested one such diagnostic in his talk, involving the alignment of the produced quarkonia. If the color octet $L = 0$ contribution dominates in $\psi$ production then, since the soft glouns are irrelevant in the decay, one expects that the produced $\psi$ should be transversally aligned. Hence the produced leptons from the decay $\psi \to \ell^+\ell^-$ should have an angular distribution proportional to $1 + \cos^2\theta$. Unfortunately, the practical situation is not so simple since about 30% of the $\psi$'s come from radiative decays of produced $\chi$'s ($\chi \to \gamma\psi$) and so this dilutes the purity of the signal. Furthermore, detecting the asymmetry in the production angle is hard experimentally for $\psi$'s produced at large transverse momentum, due to the substantial kinematical boost of the produced leptons.

Still within QCD, but now in the non-perturbative sector, we heard also of some nice work in the Symposium connected with novel quarkonia, like $B_c$ and baryons

containing two different heavy quarks $QQ'q$. If one has systems like $B_c$ or $QQ'q$ with two heavy quarks of quite different masses, then mass effects can lead to substantial differences. For instance, as Chang[8] and Oakes[9] discussed, the hyperfine splitting between $^3S_1$ and $^1S_0$ in the $B_c$ system is only about 70 MeV compared to 125 MeV in charmonium and 100 MeV in bottomonium. For the double heavy baryons, one approach discussed by Chang[8] is to consider them as bound states of a heavy diquark–light quark system:

$$QQ'q \sim \bar{3}_{QQ'}q \ .$$

This system is then not that dissimilar from a heavy-light meson, like $B_c$. However, the diquark $(bc)$ is much less tightly bound than the meson $(b\bar{c})$ [8], with

$$M_{bc} \simeq 6.6 \text{ GeV} \quad \text{versus} \quad M_{B_c \sim b\bar{c}} \simeq 6.3 \text{ GeV} \ .$$

## 3. Heavy Flavor Decays.

The physics of heavy quark systems is an important testing ground for our theoretical understanding of QCD and of the electroweak interactions. In addition, heavy quark decays offer the opportunity for exploring further the still poorly understood phenomena of CP violation. The activity in this field, which was mirrored in this meeting, roughly splits into two pieces:

i) Improvements and refinements in dynamical calculations of weak decay matrix elements by a variety of techniques: parton/quark models; chiral perturbation theory; $1/N_c$ methods; lattice calculations; and QCD sum rules.

ii) Exploration of areas where one can probe better the standard model, or look for signs of new physics. These included CP violation in charged-$B$ decays; new ways to determine the angles in the unitarity triangle; studies of non-CKM CP-violating phases; and the physics of $\tau$ lepton decays.

The talks of T. Huang[10] and W. Bardeen[11] in this Symposium provided two examples of attempts at better estimating dynamical parameters in weak decays which are of considerable phenomenological interest. Huang[10] discussed SU(3) breaking effects for the predictions of various quantities obtained by using heavy quark effective theory (HQET), using QCD sum rules as a tool. His results are as follows:

i) The ratios of weak decay constants receive about a 10% SU(3) breaking corrections

$$\frac{f_{Bs}}{f_{Bd}} = 1.18 \pm 0.05; \quad \frac{f_{Ds}}{f_D} = 1.13 \pm 0.03 \ .$$

These results are quite compatible with lattice calculations. Furthermore, as Oakes[9] pointed out, the double ratio of the above quantities is quite insensitive

to SU(3) breaking. These results are important for phenomenology since, for example, the $B_s - \bar{B}_s$ mass difference $\Delta m_s$ can be derived from the $B_d - \bar{B}_d$ mass difference and CKM parameters once $f_{B_s}/f_{B_d}$ is known.

ii) The Isgur-Wise function and the operators coefficients of the HQET Lagrangian are quite insensitive to SU(3) breaking, with corrections of order a few percent. However, Huang[10] finds that the slope parameters in the Isgur-Wise function obey $\rho_s^2 > \rho_{u,d}^2$, which is the opposite behavior of that obtained in chiral perturbation theory.

Bardeen[11] discussed another parameter of phenomenological importance for $B$ physics, the, so-called, bag constant $B_{Bd}$ which gives a measure of the $\Delta B = 2$ matrix element:

$$\langle B_d|\bar{d}\gamma_\mu(1-\gamma_5)b\bar{d}\gamma^\mu(1-\gamma_5)b|\bar{B}_d\rangle = \frac{8}{3}f_{Bd}M_{Bd}^2 B_{Bd} \ .$$

Because this matrix element enters in the expression for the $B_d - \bar{B}_d$ mass difference, changes in the value of $B_{Bd}$ affect the constraints one obtains for the CKM parameters obtained from the experimental value of this mass difference. Both lattice methods and QCD sum rules give values for $B_{Bd}$ very close to unity. Bardeen calculates this quantity using $1/N_c$ methods.

The leading contribution for $B_{Bd}$ in a large $N_c$ expansion corresponds to introducing the vacuum state in the above matrix element and leads to $B_{Bd} = 3/4$. Non-leading contributions come from the connected matrix elements involving the 2-current correlation

$$\text{corr} \ = \int d^4q\langle B_d|J_\mu(q)J^\mu(-q)|\bar{B}_d\rangle$$

To proceed, Bardeen[11] uses different techniques to evaluate the above integral in different regions of momentum $q$, matching these calculations at their interface. Writing $q^\mu = m_b v^\mu + k^\mu$, Bardeen[11] uses HQET to calculate for $\Lambda_{\text{QCD}} < k$, but uses an effective meson theory for $k < \Lambda_{\text{QCD}}$.

Both the HQET and the effective meson theory give integrals for the correction factor which are both infrared and ultraviolet sensitive and matching these contributions gives two conditions. One of them is a matching scale which turns out to be $\lambda \simeq 600\sqrt{\alpha_s}$ MeV. The other is a condition on the coupling strength in the effective theory and Bardeen obtains $g^2 = 1/3$. Remarkably, because of this second matching condition, the result for $B_{Bd}$ that Bardeen[11] obtains is unaffected by the nonleading corrections in $1/N_c$:

$$B_{Bd} = \frac{3}{4}[1 - 0.1(1-3g^2)] \longrightarrow \frac{3}{4} \ .$$

As Bardeen points out, it is not clear how general this result is. For instance, in his effective meson calculation he has included $B_d^*$ states but not, for instance, $B_d^{**}$ states.

The inclusion of these further states could change the coupling strength matching condition and thus the result for $B_{Bd}$. Nevertheless, it is troubling that there appears to be a discrepancy between the value obtained for $B_{Bd}$ in lattice and QCD sum rules calculations and in this $1/N_c$ calculation.

In the Symposium Lam[12] also discussed the large $N_c$ limit, but applied to baryons which in this limit are just large collections of quarks: $B \sim N_c q$. As $N_c \to \infty$ these states are necessarily heavy, if the quarks carry any mass. Lam described in particular how to reconcile, in a special kinematical limit, the fact that baryonic decays to $n$ mesons are highly suppressed in the large $N_c$ limit, with

$$A(B \to B' nM) \sim O\left(N_c^{\frac{2-n}{2}}\right) ,$$

while individual Feynamn graphs are all of $O(N_c^{n/2})$ and, apparently, grow with $N_c$. The reconciliation is effected by having an infinite tower of resonances in the theory in the large $N_c$ limit, with all the $MBB^*$ couplings being appropriately related.

Also somewhat theoretical was the nice discussion of C.-S. Huang[13] of how to recover the results of HQET in a Bethe-Salpeter formalism. One expects this to emerge in an analogous way that one recovers in the non-relativistic limit the Schrödinger equation from the Bethe-Salpeter equation. Nevertheless, it was nice to see how this obtains in detail, recovering both the spin symmetry as $M_Q \to \infty$ (provided one has vector or scalar kernels) and the HQET form of the $1/M_q$ corrections.

Huang[13] applied this covariant formalism to a model calculation of exclusive semileptonic decays, where he extracted the Isgur-Wise function, and to other heavy quark non-leptonic decays, like $D^* \to D\pi$. Similar calculation to these were discussed at the Symposium by C.-S. Kim[14], who used a parton model for his calculations, and by L.-H. Chan[15] who used an effective low-energy Lagrangian similar to that discussed by Bardeen[11].

Kamal[16] also presented a model investigation, in his case concerning the color suppressed decays of the $B$ mesons into $\psi K$ and $\psi K^*$. Kamal remarked that the usual calculation, where one drops the color pieces in the effective Lagrangian after Fierzing the currents and where one uses factorization, cannot reproduce the experimental values for either the ratio of these modes or the polarization in the $\psi K^*$ mode:

$$R = \frac{BR(B \to \psi K^*)}{BR(B \to \psi K)} = 1.71 \pm 0.34; \qquad P_L(B \to \psi K^*) = 0.78 \pm 0.07 .$$

These two assumptions (using $N_c = 3$) give a small $a_2$ amplitude, with

$$a_2 = c_2 + \frac{1}{N_c} c_1 \simeq 0.1$$

What Kamal[16] pointed out was that everything works out–both here and in color suppressed $D$-decays–if there is about a 10% non-factorizable contribution and an

analogous $O(10\%)$ contribution from the color pieces in the effective Lagrangian. These contributions, effectively, conspire to change the $a_2$ amplitude to a new effective amplitude, with

$$a_2^{\rm eff} \simeq c_2 \ .$$

So Kamal's results are similar to just imagining dropping the $1/N_c$ contributions–a suggestion made earlier in the literature[17].

Much more model-independent was the discussion of Paschos[18] at the Symposium of inclusive semileptonic $B$-decays. Because one is summing over all hadronic final states, the inclusive rate can be written in terms of a current commutator taken between $B$ states:

$$W_{\mu\nu} = \int d^4x e^{-iqx} \langle B|[J_\mu(x), J_\nu(0)]|B\rangle \ ,$$

where $q^\mu$ is the momentum transfer to the final lepton pair. This quantity can be calculated in a controlled way for most of the allowed phase space by using a combination of a light-cone expansion and HQET. Thus, one expects that the inclusive semileptonic rate should be reliably calculable in terms of the parton model, augmented by the matrix elements of $O(1/m_b)$ operators [$D^2$ and $\sigma \cdot G$] arising from the light-cone expansion. Unfortunately, these expectations are not realized in practice since the experimental semileptonic branching ratio

$$B_{sL} = \frac{\Gamma(B \to X\ell\nu_e)}{\Gamma(B \to \ \text{all})} = 10.6 \pm 0.3$$

is quite a bit smaller than the theoretical prediction of 12-13%.

Paschos[18] discussed some possibilities for reconciling theory with experiment. This can happen readily if one, somehow, underestimated the strength of the non-leptonic $B$-decays. The favored idea here is that the mode $b \to c\bar{c}s$ is underestimated. However, to bring theory and experiment in concordance one would need to boost up this mode so much that it would lead to too much charm production ($N_c \sim 1.3$), in conflict with observation. It is possible that the discrepancy is the efffect of new physics, where a favored effective operator is that given by

$$L_{\rm eff} \sim \frac{1}{M^2_{\rm new}} (\bar{b}s)_R(\bar{q}q) \ .$$

However, it may also just be that we, again, have failed to correctly calculate the relevant non-leptonic matrix element. History perhaps gives credence to this last, more humble, hypothesis. For kaons, the $\Delta I = 1/2$ enhancement is a factor of 20 which, even today, is only partially understood. We also have not really totally explained the factor of 2 difference between the charm lifetimes, $\tau(D^+)/\tau(D_0) \sim 2$. So perhaps we should not be too concerned by a 20% discrepancy in the semileptonic $B$-decays!

D.-S. Du[19] in his talk at the Symposium suggested that one should consider anew the possibility of having rather large CP-violating asymmetries in charged $B$-decays.

This is an old suggestion[20] which, however, seems to be difficult to realize in practice. To obtain a CP-violating asymmetry in $B^\pm$-decays requires the interference of two amplitudes with both **different** weak CP-violating phases and strong rescattering phases. Although this occurs in practice, in general one of the amplitudes or one of the phase differences is small and the net asymmetry is then also small. Du[19] suggests that this may not happen for decays like $B^\pm \to \pi^\pm \pi^o$ where one is interfering a spectator decay amplitude with a (space-like) Penguin amplitude. Du gets a large effect by assuming that the size of the space-like Penguin amplitude is related to the Brodsky-Lepage[21] form factor:

$$\langle \pi\pi | J | 0 \rangle \sim \frac{i\alpha_s}{M_B^2} \ .$$

This gives him an amplitude which is comparable in size to the spectator decay amplitude and in which the rescattering phase is maximum. Because the two amplitudes in question involve $V_{ub}$ and $V_{td}$, respectively, the weak phases are also comparable. So, in principle, one could get large effects. Unfortunately, it is difficult to judge how reliable the Penguin estimate of Du[19] is. At any rate, he has raised an interesting issue.

Tau decays were also discussed at the Symposium, both as a beautiful laboratory for applying current algebra and dispersion relation techniques[22] and as a place to look for new physics[23]. Truong[22] emphasized that the current algebra soft pion relation in the limit of $p^\mu \to 0$:

$$\langle B\pi | V_\mu | A \rangle = \frac{1}{f\pi} \langle B | A_\mu | A \rangle \ ,$$

when used with the Padé techniques to build-in unitarity, can be very powerful. Indeed, by these means it is possible to make successful predictions for multipion $\tau$-decays ($\tau \to n\pi\nu_\tau$), including resonance channels, like $\tau \to \pi\rho\nu_\tau$. Nelson[23] instead concentrated on what limits on new physics could be obtained from $\tau$-decays at the proposed Beijing tau-charm factory. He showed that, by looking at the $\tau \to \rho\nu_\tau$ and $\tau \to A_1\nu_\tau$ decays and analyzing the $\rho$ and $A_1$ polarization through their further decays, one can obtain limits on the scale associated with new V-A interactions of the $\tau$ which are of $O(\Lambda \sim 1 \text{ TeV})$. Nelson[23] also showed that one could test for possible CP-violating asymmetries in the charged $\tau$-decays to quite a reasonable level. For instance, writing the amplitude for $\tau^\pm \to \rho^\pm \nu_\tau$ as $r^\pm = |r| e^{i\phi}$, at a tau-charm factory one could hope to determine $\delta r/r$ to about 0.1% and $\delta\phi$ to about $1^\circ$.

## 4. Electroweak Symmetry Breaking.

The third subject of great interest at the Symposium was electroweak physics. Here there are a few facts which were agreed by all the speakers, either implicitly or explicitly:

i) The standard model gives an amamzingly accurate description of a large body of precise electroweak data[24,25]. An example being provided by the very accurate value of $sin^2\theta_{\rm eff} = 0.2315 \pm 0.0004$.

ii) The physics underlying the breakdown of $SU(2) \times U(1) \rightarrow U(1)_{\rm em}$ occurs at scale of $O(1$ TeV$)$.

iii) The large mass of the top quark, with $m_t \sim O(v)$ and where $v = (\sqrt{2}\, G_F)^{-1/2} \simeq 250$ GeV is the scale associated with the Higgs vev, is significant. Although what exactly this is telling us is not yet totally clear[26,27].

The focus of the discussion at the Symposium was on the **disputable aspects** of the above points. For example, are there hints of small discrepancies with the standard model in the data? Or, what really is the physics which is at the root of the symmetry breakdown? Or, what is the real significance of having top so heavy?

Probably the central issue of particle physics today is what is the mechanism which causes the $SU(2) \times U(1)$ breakdown. Two camps exist. Partisans of the first camp believe that the breakdown is due to the vev of some elementary scalar(s) field(s).[25] This is the original mechanism suggested for the spontaneous breakdown of the standard model. However, to make this mechanism natural the belief now is that one needs to have also some supersymmetry which survives to low energy. Partisans of the second camp believe instead that the spontaneous breakdown of $SU(2) \times U(1)$ is due to the formation of condensates of some underlying fermions[28]. That is, the breakdown of $SU(2) \times U(1)$ is dynamical. It is possible that what condenses to break the symmetry is just $\langle t\bar{t} \rangle$, but generally it is assumed that the condensing fermions are fermions of a new theory–technicolor.

If the first option above is the truth and one has some low energy supersymmetry, then eventually one should see plenty of signals. All known excitations will have superpartners and their spectrum will inform us of how precisely the supersymmetry is broken down in nature. Furthermore, since to implement the supersymmetry one needs at least 2 Higgs doublets, one should also observe the scalar excitations connected with an extended Higgs sector.[25] In general, a relatively light Higgs boson ($M_h \leq M_Z$) is symptomatic of supersymmetry. One knows from direct searches at LEP that the standard model Higgs boson has a mass $M_H > 65$ GeV. As Ellis[25] discussed at the Symposium, from indirect fits to precision electroweak data one infers that $M_H = 76^{+100}_{-50}$ GeV. Optimistically, he concluded that such a "light Higgs" perhaps is already a hint of supersymmetry. Whether this is so only time (and more data!) will tell.

The breakdown of the electroweak symmetry by a Higgs vev which is stabilized by supersymmetry is, in many respects, a much "safer" option than dynamical symmetry breaking. Principally this is because it does not tie the scale of $SU(2) \times U(1)$ breaking to the physics scale responsible for generating the Yukawa couplings of the Higgs to

the fermions, which are responsible for fermion masses. This cannot be avoided when the symmetry breaking is dynamical and, in these latter theories, one is forced to have the fermion mass generation scale near to the $O(\text{TeV})$ scale of $SU(2) \times U(1)$ breaking.

Simmons[28] discussed at the Symposium how the large mass of the top makes life even more difficult. Typically, when the electroweak breakdown is caused dynamically, one generates fermion masses though effective 4-fermion interactions between the ordinary quark and leptons and a new set of fermions (technifermions) whose condensation causes the breakdown. This ETC mechanism[29] provides an effective Lagrangian of the form

$$L_{\text{eff}} \sim \frac{1}{M^2}(\bar{T}T)(\bar{\psi}\psi)$$

where $M$ is the scale of the ETC interactions which connect the ordinary fermions $\psi$ with the technifermions $T$. The breakdown of $SU(2) \times U(1)$ occurs as a result of the formation of a $\langle\bar{T}T\rangle$ condensate. Because of the above effective interactions, these condensates also give mass to the ordinary fermions. Since top has such a large mass and

$$m_t \sim \frac{\langle\bar{T}T\rangle}{M^2} \ ,$$

the fermion mass generating scale $M$ cannot be very large.[28] Because the electroweak breaking scale associated with the $\langle\bar{T}T\rangle$ condensate is of $O(\text{TeV})$ [i.e. $\langle\bar{T}T\rangle \sim (\text{TeV})^3$] the scale $M \sim O(10 \text{ TeV})$, at most. The presence of such "low scales" for new physics associated with fermion mass generation, in general, produces unwanted flavor changing neutral currents and one must devise rather clever schemes[30] to avoid these troubles. Furthermore, the technicolor condensates themselves produce small changes in the expectations of precision electroweak tests and these changes are not favored experimentally. For instance, as Kang[24] discussed, the so-called $S$ parameter is, in general, positive as a result of having $\langle\bar{T}T\rangle$ condensates, while data prefers $S < 0$.

Simmons[28] pointed out an especially serious problem for classes of ETC models precisely in the area where there appears to be some discrepancy between the data and the standard model.[24,25] This is in the ratios of the widths of the $Z$ into $b\bar{b}$ and $c\bar{c}$ states to the total width. Experimentally, one has

$$R_b = \frac{\Gamma(Z \to b\bar{b})}{\Gamma(Z \to \text{ hadrons})} = 0.222 \pm 0.002; \qquad R_c = \frac{\Gamma(Z \to c\bar{c})}{\Gamma(Z \to \text{ hadrons})} = 0.154 \pm 0.008,$$

while the standard model expectations are centered around 0.216 and 0.173, respectively. Simmons noted that for models where the ETC interactions commute with $SU(2)$, then the same interactions which give a mass to the top also give a specific shift to $R_b$, but no shift in $R_c$. Unfortunately, these models give a shift (of about 4%) in the **wrong** direction and therefore are strongly disfavored by the data. One can, however, invent models where the ETC interactions and $SU(2)$ do not commute and

change the sign of the $R_b$ shift [essentially, one needs to change a $\vec{\tau}\cdot\vec{\tau}$ interaction to a $1\cdot 1$ interaction]. However, the resulting models are a bit recondite in that different families are treated differently and one may run into some problems with universality.

## 5. Physics of Top.

The discovery of the top quark at Fermilab[31] was one of the year's highlights. The results of CDF and DO are as follows:

$$\begin{array}{lll} m_t = 178 \pm 11 \pm 9 \text{ GeV}; & \sigma_{t\bar{t}}(m_t) = 6.8^{+3.6}_{-2.4} \text{ pb} & \text{(CDF)} \\ m_t = 199^{+19+14}_{-21-21} \text{ GeV}; & \sigma_{t\bar{t}}(m_t) = 6.4 \pm 2.2 \text{ pb} & \text{(DO)} \end{array}$$

At the Symposium the sensitivity of these results to possible new physics contributions were discussed by Parke[26] and C.-S. Li[32], who specifically considered the effects of possible supersymmetric corrections to the top production cross section. B.-L. Young[27] instead speculated on possible non-standard couplings for the top, which may be more evident because of its large mass.

Although speculation of new physics associated with the top is fair game, there is already really not too much room to manuever. For instance, the combined value for the top mass coming from the CDF and DO measurements, $\langle m_t\rangle = (181\pm 12)$ GeV is actually in quite good agreement with that obtained through the precision electroweak tests (when the Higgs mass is considered a free parameter) reported by Ellis[25]: $m_t = (155 \pm 14)$ GeV. The average of both these values gives a top mass of $\langle\langle m_t\rangle\rangle = 172 \pm 10$ GeV. For this mass the latest calculation of the top cross section reported by Berger[2] here, of $\sigma_{t\bar{t}}(m_t) = 5.5 \pm 0.3$ pb, is in reasonable agreement with the CDF and DO values. So, it could well be that also for top everything is standard!

The discussion of Parke[26] at the Symposium emphasized what physics could explain possible disagreements between theory and experiment. Although he presented a more speculative interpretation for the present data, this exercise is very useful nevertheless. As usual, a good way to test sensitivity to new physics is to introduce contact terms describing new interactions of the top with the ordinary quarks, respecting the symmetries of the standard model. Parke[26] discussed 4-fermion interactions of the type

$$L_{\rm eff} = \frac{g_3^2}{\Lambda_1^2}(\bar{q}1q)(\bar{t}1t) + \frac{g_3^2}{\Lambda_8^2}(\bar{q}\lambda q)(\bar{t}\lambda t)$$

and indicated that present data bounds the scales $\Lambda_1$ and $\Lambda_8$ to be above a TeV. He also discussed more specific models, like the coloron model[33] where the color $SU(3)$ group of QCD arises as a result of the spontaneous breakdown of an $SU(3)\times SU(3)$ group. The octet of gauge bosons which acquire mass–the colorons– have a mass $M \sim \Lambda_8$ but have different couplings to ordinary quarks ($\sim \tan\theta$) than to top ($\sim \cot\theta$). Such a coloron model[33] predicts distinctive transverse momentum distortions for top production and structure in the invariant mass of the produced $t\bar{t}$-pairs.[26]

B.-L. Young[27] discussed another aspect of possible anomalies connected with top. If the symmetry breakdown of the electroweak theory is dynamical, it is natural to expect anomalous interactions of the Nambu-Goldstone bosons with the fermions in the theory

$$L_{\text{eff}} \sim \frac{\kappa}{\Lambda}\bar{\psi}\gamma_\mu\psi\partial^\mu\xi + \ldots$$

By the equivalence theorem, discussed here by Y.-P. Kuang[34], these couplings eventually give rise to anomalous vertices of the fermions with the gauge bosons. For top these anomalous vertices could be of significant strength, since one expects $\kappa \sim O(m_t/\Lambda)$ and $\Lambda$ to be in the TeV range. Therefore, because of the large mass of top, one could be sensitive to new phenomena connected with the way the electroweak symmetry breaks down. These anomalous vertices, as Young[27] discussed, could be responsible for the small discrepancy in $R_b$ and could also give rise to other phenomena, like flavor changing decays of the top, which may be observable some day.[35]

## 6. Concluding Remarks.

My conclusions are very simple. This has been an exciting and fun meeting to be at, with plenty of physics bubbling up! Such a meeting would not have been possible without all the hard work done by the Organizers. On behalf of all of the participants, I would like to thank them for their splendid hospitality.

## 7. Acknowledgments.

This work was supported in part by the Department of Energy under Grant No. FG03-91ER40662. I would like also to thank Tuo Huang for having made my brief stay in Beijing so enjoyable.

## References.

1. S. J. Brodsky, in these Proceedings.
2. E. Berger, in these Proceedings.
3. M. B. Wise, in these Proceedings.
4. E. Laenen, J. Smith and W. L. van Neerven, Nucl. Phys. B**369** (1992) 543; Phys. Lett. **321B** (1994) 254.
5. E. Berger and H. Contopanagos, Phys. Lett. B361 (1995) 115.
6. CDF Collaboration, F. Abe *et al.*, Phys. Rev. Lett. **69** (1993) 3704; **71** (1993) 2537.
7. G. T. Bodwin, E. Braaten, T. C. Yuan and G. P. Lepage, Phys. Rev. D**46** (1992) R3703; G. T. Bodwin, E. Braaten and G. P. Lepage, Phys. Rev. D**51**

(1992) 1125; E. Braaten and S. Fleming, Phys. Rev. Lett. **74** (1995) 3327; S. Fleming and I. Maksymyk, MADPH-995-922.

8. C.-H. Chang, in these Proceedings.
9. R. J. Oakes, in these Proceedings.
10. T. Huang, in these Proceedings.
11. W. A. Bardeen, in these Proceedings.
12. C.-S. Lam, in these Proceedings.
13. C.-S. Huang, in these Proceedings.
14. C. S. Kim, in these Proceedings.
15. L.-H. Chan, in these Proceedings.
16. A. N. Kamal, in these Proceedings.
17. A. J. Buras, J.-M. Gerard and R. Rückl, Nucl. Phys. B**268** (1936) 16.
18. E. A. Paschos, in these Proceedings.
19. D.-S. Du, in these Proceedings.
20. L. L. Chau and H. Y. Cheng, Phys. Rev. Lett. **59** (1987) 958.
21. G. P. Lepage and S. J. Brodsky, Phys. Lett. **87B** (1979) 359.
22. T. N. Truong, in these Proceedings.
23. C. A. Nelson, in these Proceedings.
24. K. Kang, in these Proceedings.
25. J. Ellis, in these Proceedings.
26. S. Parke, in these Proceedings.
27. B.-L. Young, in these Proceedings.
28. E. Simmons, in these Proceedings.
29. S. Dimopoulos and L. Susskind, Nucl. Phys. B**155** (1979) 237; E. Eichten and K. Lane, Phys. Lett. B**90** (1980) 125.
30. See, for example, B. Holdom, in the Proceedings of the 1991 Nagoya Spring School on Dynamical Symmetry Breaking, ed. K. Yamawaki (World Scientific, Singapore, 1992).
31. CDF Collaboration, F. Abe *et al.*, Phys. Rev. Lett. **74** (1995), 2626; DO Collaboration, S. Abachi *et al.*, Phys. Rev. Lett. **74** (1995) 2632.
32. C.-S. Li, in these Proceedings.
33. C. T. Hill and S. J. Parke, Phys. Rev. D**49** (1995) 4454.
34. Y.-P. Kuang, in these Proceedings.
35. T. Han, R. D. Peccei and X. Zhang, Nucl. Phys. B**454** (1995) 527.